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THE INTERACTION OF INTENSE LASER BEAMS WITH FREE ELECTRONS

by Edward S. Sarachik

Electronics Research Center

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ABSTRACT

The field of high intensity electrodynamics is introduced and surveyed. Various aspects of the interactions of intense laser beams with free electrons are considered, including harmonic generation, refraction, the Kapitza-Dirac effect, and quantum phenomena. A complete bibliography is included.

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SUMMARY

In this report the field of high-intensity electrodynamics is critically introduced and surveyed by considering the interactions of free electrons with very intense laser beams. The concept of high intensity is developed by noting that the corrections to ordinary Compton scattering depends on an intensity parameter, called q^2. To lowest order in q^2, harmonic generation from free electrons, intensity-dependent frequency shifts in Compton scattering, and the reflection and refraction of electrons from laser beams are considered. Effects of a purely quantum mechanical nature, which are influenced by the presence of an intense laser beam, are also discussed. Finally, the Kapitza-Dirac effect, an example of stimulated Compton scattering, is discussed.

A bibliography of laser-free electron interactions--complete to December 1968--is included in the report.

I. INTRODUCTION

The interaction of photons with electrons has been the focus of much of the fundamental physics research of this century. By the early 1900's, the scattering of low-energy photons (100 keV and below, including the visible at 1 eV) was well accounted for by the classical wave theory of Thomson scattering wherein an electron was driven to oscillate at the frequency of the incident wave and consequently reradiated at this same frequency. In the 1920's, deviations from Thomson scattering above 100 keV were accounted for by the photon picture which treated the scattering as an inelastic collision of a single photon with a single electron, the photon being downshifted in frequency as a result of the collision. This billiard ball model, called Compton scattering, was further refined by Dirac's theory of the electron in the 1930's and ultimately, in the 1940's, by the development of covariant perturbation theory resulting in a complete and accurate theory of quantum electrodynamics. At this time, it is true that up to the highest energies presently obtainable, the interaction of individual photons and electrons is completely understood and predictable to arbitrary precision.

In developing this exceedingly powerful and accurate theory, no account had to be taken of the coherence aspect of the photons since existing photon sources produced incoherent photons one at a time. A theory describing a single photon interacting with a single electron is perfectly adequate under these conditions. It was not until the development of lasers that the properties of light beams became important, and it was not until this year, with the development of picosecond pulses that light beams of high intensity were actually achieved. Concurrent with the development of high-power lasers during the last few years, a few theorists became interested in laser-electron interactions and a small body of literature became available. It is the purpose of this report to review what is known about the theory and to point out some of the surprising things it predicts. It should be borne in mind that, at this writing, not a single experiment in this entire field of high-intensity electrodynamics has yet been performed. Thus, the theories must be considered tentative and, as higher power lasers become available and more experiments are performed, may be subject to massive revision.

II. LINEAR THOMSON SCATTERING

The theory of ordinary Thomson scattering (ref. 1) assumes that a free electron at rest in a linearly polarized monochromatic plane wave of frequency ω responds only to the electric field

$$\vec{E} = \vec{E}_o \cos \omega t \ .$$

The force on the electron is therefore $e\vec{E}$ and the electron consequently executes simple harmonic motion at frequency ω along the $\vec{E}_o$ direction:

$$\vec{x} = -\frac{1}{\omega^2} \frac{e}{m} \vec{E}_o \cos \omega t \ .$$

The radiation from such an oscillating charge can be simply calculated by the classical Larmor radiation formula (ref. 1), and the calculation yields the differential cross section (final polarization not measured):

$$\frac{d\sigma_{\text{Thomson}}}{d\Omega} = r_0^2\left(\sin^2 \phi + \cos^2 \theta \cos^2 \phi\right)$$

where $r_o = e^2/(mc^2)$ is the classical electron radius, ϕ is the azimuthal angle with respect to the incident polarization, and θ is the scattering angle. The geometry is illustrated in Figure 1. The total Thomson cross section is

$$\sigma_{Thomson} = \frac{8\pi}{3} r_0^2 \cong 10^{-25} \text{ cm}^2 .$$

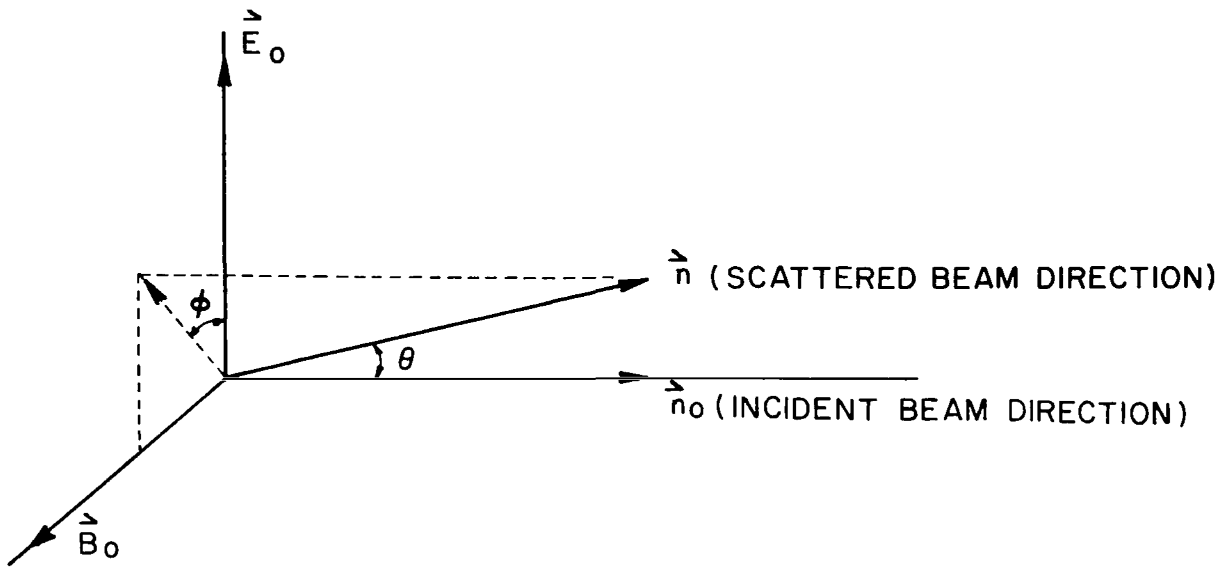

Figure 1

The basic assumption in this discussion has been that the $(e/c)(\vec{v} \times \vec{B})$ force due to the magnetic field could be neglected; i.e., that

$$\frac{ev|B_o|}{c} \ll e|E_o| ,$$

and since $|E_o| = |B_o|$ in these units, that $v/c \ll 1$. Thus, the basic assumption is that the problem be completely non-relativistic. Since

$$v_{max} \sim \frac{e|E_o|}{m\omega}$$

this assumption will hold true as long as the quantity

$$q = \left[\frac{e^2 I \lambda^2}{\pi m^2 c^5}\right]^{1/2} \sqrt{2} \ll 1 ,$$

where I is the beam intensity $[I = (cE_0^2)/(8\pi)]$ in W/cm^2. It can be seen that the condition that the laser-electron interaction be describable by linear Thomson theory is that the intensity be so low that $q^2 \ll 1$. High intensities are defined as that region of intensities for which q^2 is of order unity. As an illustration, note that an ordinary light bulb corresponds to $q^2 \approx 10^{-18}$. Much more will be said about this crucial parameter q^2 in the following sections.

III. NON-LINEAR THOMSON SCATTERING — SECOND HARMONIC GENERATION

When the intensity of the beam increases to a point that the $q^2 \ll 1$ condition of the previous section no longer holds, the situation changes drastically. The electron now becomes relativistic, the $(e/c)(\vec{v} \times \vec{B})$ force can no longer be neglected, and the problem, in fact, becomes non-linear. The Lorentz force $\vec{F} = e\vec{E} + (e/c)(\vec{v} \times \vec{B})$ on the electron now contains not only the term linear in $\vec{E}$, but also a term at least quadratic in the field strengths due to the implicit dependence of $\vec{v}$ on $\vec{E}$. At still higher intensities, $\vec{v}$ is a more complicated function of $\vec{E}$ and $\vec{B}$ and the problem becomes even more non-linear.

If one is confined to the lowest order of non-linearity in which $\vec{v} \propto \vec{E}$, it can be easily seen that the electron becomes subjected to a magnetic force $\vec{F} \propto \vec{E} \times \vec{B}$ and will, therefore, have an oscillatory component along the beam direction at twice the incident radiation frequency. This oscillating electron will then emit radiation at a frequency twice the incident frequency, and the interesting phenomenon of second harmonic generation (SHG) off free electrons occurs.

To be more precise, the electron is driven into oscillation with frequency ω along the $\vec{E}_o$ direction by the electric force and with frequency 2ω along the beam direction by the magnetic force. The combined motion is a "figure 8," the equation of which can be written (refs. 2-4):

$$\vec{x} = \frac{1}{k_0}\left[-\frac{eE_0}{mk_0'}\left(\cos k_0'\tau\right)\vec{e}_0 - \frac{1}{8}\left(\frac{k_0}{k_0'}\right)^2 q^2\left(\sin 2k_0'\tau\right)\vec{n}_0\right]$$

where

$$\vec{e}_o = \frac{\vec{E}_o}{|E_o|} = \text{the initial polarization direction,}$$

$\vec{n}_o$ = the beam direction, and

$$k_o' = k_o\sqrt{1+\frac{1}{2}q^2}\ ,\ k_o = \frac{2\pi}{\lambda} = \frac{\omega}{c}\ ,\ \text{and}\quad q^2 = \frac{2e^2I\lambda^2}{\pi m^2c^5}\ ,\ \text{and}$$

τ = the proper time of the electron, the motion taking place in a frame in which the electron is on the average at rest.

It can be seen that the second harmonic motion becomes large only when q^2 is of order unity. This is precisely the condition that the theory of linear Thomson scattering is no longer adequate, as was seen in Section II.

The radiation pattern due to a charge going in a figure 8 has been calculated (ref. 4) and results in the following differential cross-section (final polarization not measured) for the second harmonic to order q^2:

$$\frac{d\sigma}{d\Omega}(2\omega) = r_0^2q^2\,\frac{\sin^2\theta}{4}\Big[1 + 16\cos^4\phi - 8\cos^2\phi\cos\theta$$
$$- 16\cos^4\phi + 16\cos^4\phi\cos^2\theta\Big]$$

and a total cross-section

$$\sigma(2\omega) = \frac{14\pi}{5}\,r_0^2q^2\ .$$

If the detector is placed along the direction of the incident polarization ($\phi = 0$), the differential cross-section can be plotted as a function of the scattering angle θ as shown in Figure 2.

It should be noted that the cross-section becomes comparable to the linear Thomson cross-section only for $q^2 \sim 1$. It may also

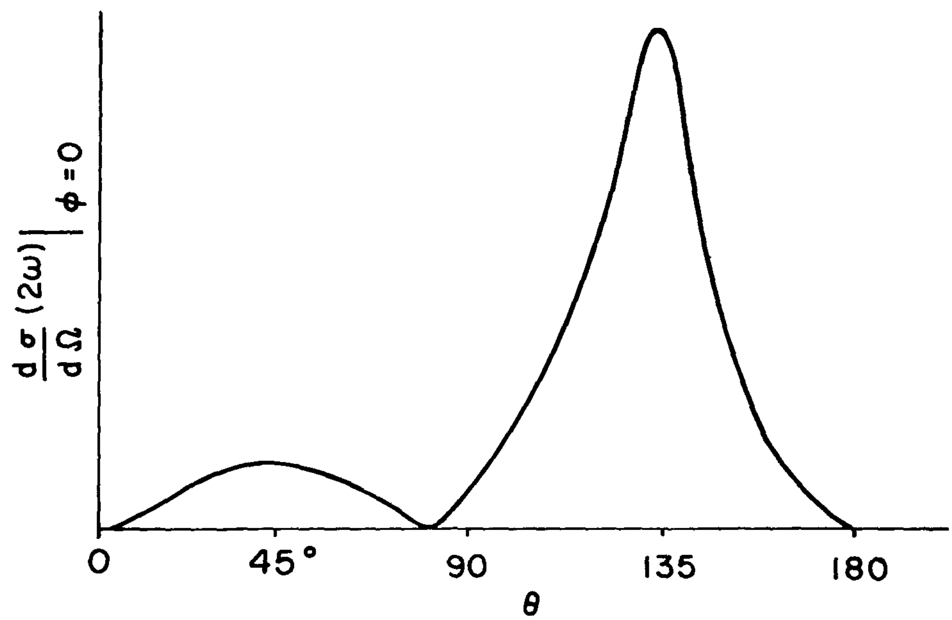

Figure 2

be noted that the differential cross-section vanishes in the forward and backward directions and has a large backward peak at 130 degrees. The forward-backward vanishing is characteristic of a radiating dipole induced along the beam direction and is understandable, but the prediction of the backward peak is surprising.

This classical calculation has only been carried out to order q^2 in a perturbation expansion assuming q^2 of order unity but less than 1. At q^2 near unity, higher order terms have to be included, and at $q^2 = 1$, the perturbation expansion breaks down and the above results are no longer valid.

The calculation has also been performed quantum mechanically (refs. 5-7) by scattering a photon off an electron, the wave function of which is chosen to be that of an electron in an intense monochromatic wave, the Volkov wave functions (refs. 8,9). The solution must again be expanded in a perturbation series in q^2 and, to lowest order in q^2, results identical to the classical calculations were obtained. The spin of the electron does not seem to matter to order q^2, although it may enter in higher order.

Many papers have been written expanding on the basic calculations outlined above both classically (refs. 10, 11, 12-22) and quantum mechanically (refs. 14, 23-28) and on various other facets of non-linear Thomson scattering, especially the interesting question of radiation reaction (refs. 12, 13, 15, 29-31).

IV. THE PARAMETER q^2

It was shown that second harmonic generation off free electrons became sizable only when q^2 grew to order unity, where q^2 was defined as

$$q^2 = \frac{2e^2 I\lambda^2}{\pi m^2 c^5} . \tag{1}$$

Numerically, $q^2 = 3\times10^{-11}\, I\lambda^2$, with the beam intensity I in W/cm^2 and the wavelength λ in cm. For $\lambda = 10{,}000$ Å $= 10^{-4}$ cm, an intensity of 3×10^{18} W/cm^2 is necessary for $q^2 = 1$. This intensity is within the reach of present-day focused mode-locked lasers.

The parameter q^2 is a measure of whether or not the region of high-intensity (non-linear) electrodynamics has been entered. It can be seen immediately that, since $q^2 \propto \lambda^2 \propto 1/\omega^2$, in high energy photon physics q^2 is very small. It becomes large only for high intensity, low energy electromagnetic waves. It can be seen very clearly why the field of high intensity electrodynamics had to await the development of high intensity lasers rather than high energy accelerators. In this low energy region, only microwaves and light can be produced with any reasonable intensity and, at this moment, only light can reach $q^2 = 1$.

More insight into the meaning of q^2 can be gained by considering the following forms alternate to the form (1) above:

$$q^2 = \frac{2r_o\lambda_c\lambda\rho}{\pi} \tag{2}$$

where

ρ = photon number density

$\lambda_c = h/mc$ = electron Compton wavelength

$r_0 = e^2/mc^2$ = classical electron radius.

and

$$q^2 = \frac{2}{\pi} \frac{\mathscr{E} r_0 \lambda^2}{mc^2} \tag{3}$$

where

$\mathscr{E}$ = energy density of the wave.

In form (2), q^2 is the number of photons in a volume $r_0 \lambda_c \lambda$, and in form (3) it becomes the ratio of the electromagnetic energy continued in a volume $r_0 \lambda^2$ to net mass energy of an electron. This last form is the most revealing for it states that when the electromagnetic energy in a volume characteristic of the interaction of a wave field with an electron becomes comparable to the rest energy of the electron, the region of high intensity has been entered. Note that in this high intensity region the radiation reaction on the electron may still be neglected since the radiation reaction only becomes large when the electromagnetic energy in a volume occupied by the electron, r_0^3, becomes comparable to the rest mass. Thus, for

$$q^2 \ll \frac{\lambda_c \lambda}{r_0^2} = 10^{12} \ ,$$

the radiation reaction may be neglected.

Finally, it may be noted that at really low frequencies, q^2 seems to diverge. But it is to be expected that as the frequency decreases, the energy inside a volume λ^3 should be proportional to the total energy of the field. Thus, for large λ, $\mathscr{E} \propto 1/\lambda^3$, so that q^2 now goes to zero as may be seen from form (3). Thus, although 60 cycle waves (λ = 2500 km) may seem to be the ideal choice, the formalism is in fact no longer valid at this frequency, and microwaves seem to be the smallest frequencies practically capable of producing high intensity beams.

V. NON-LINEAR THOMSON SCATTERING -- HIGHER HARMONICS

If the perturbation expansion is carried out to higher order in q^2, the cross-sections for the production of higher harmonics off free electrons can be estimated (refs. 5, 6). The cross-section for the nth harmonic goes as

$$\sigma(n\omega) \approx r_0^2 (q^2)^{n-1} ,$$

a result which obviously must break down at $q^2 = 1$; at $q^2 = 1$, all higher harmonics cannot be copiously produced because there is only a finite amount of energy in the incident beam. This is the usual situation involving the breakdown of a perturbation expansion when the expansion parameter becomes unity. With a highly simplified model, the prediction has been made (ref. 6) that the most probable harmonic is the $(q^2)^{3/2}$ harmonic, but the distribution of energy among the harmonics is not understood.

It is fair to say at this time that due to the breakdown of perturbation theory there is no theoretical understanding -- either classically or quantum mechanically -- of high intensity electrodynamics for $q^2 \geq 1$. The best place to try to understand this area theoretically would be the classical electron interacting with a classical electromagnetic field since the situation is much more clearly defined than in the quantum case. Experiments in the $q^2 \geq 1$ area would be of great use and should become performable within the next few years.

VI. INTENSITY-DEPENDENT FREQUENCY SHIFT

It has been shown that the scattering of an intense laser beam off free electrons gives rise to linear Thomson scattering and to scattering at the harmonics of the incident beam. The linear Thomson effect should predict that the scattered wave has a strong component at precisely the incident frequency. But if one looks a bit more closely, he can see that the frequency is, in fact, shifted an amount proportional to q^2.

When the equations of motion of an electron in a monochromatic plane wave are solved, it is found that the electron, initially at rest in the lab frame, picks up a component of velocity in the beam direction (ref. 32). This velocity -- of magnitude $\overline{v} = q^2 c/2$ -- is due to the force extended on the electron while the beam is being turned on and is constant in a beam of constant intensity. This will be discussed in the next section. Since the electron is now moving (Figure 3) with respect to the lab frame, any radiation it emits will be Doppler-shifted by an amount

$$\nu' = \frac{\nu}{1 + \frac{\overline{v}}{c} \cos \theta} = \frac{\nu}{1 + \frac{q^2}{2} \cos \theta}$$

so that for $0 \ll q^2 \ll 1$,

$$\frac{\Delta\nu}{\nu} = \frac{1}{2} q^2 \cos\theta$$

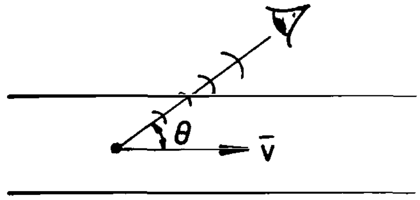

Figure 3

where θ is the scattering angle.

From form (1) of the previous section, it can be seen that q^2 depends on the intensity directly so that this frequency shift is, in fact, an intensity-dependent frequency shift (IDFS).

Including this additional motion of the electron will not affect the discussion of second harmonic generation given in section III because the transformation from the frame of the moving electron to the lab frame corrects the cross-sections by terms of order $\bar{v}/c = q^2/2$. Since the second harmonic cross-section is already of order q^2, any such corrections are of higher order in q^2. Thus, in calculating the cross-sections for the second harmonic and all higher harmonics, this additional motion does not have to be included. It does have to be included when calculating the lowest order corrections to scattering at the fundamental frequency since the cross-section itself is zero-th order in q^2. All the higher harmonic frequencies will also be shifted by this additional beam direction electron motion.

Note also that the existence of the IDFS has been challenged (refs. 27, 33) and many theoretical papers have been written on both sides of the question (refs. 32-45). The basis of this challenge is quantum mechanical and involves the properties of the electron wave function in a monochromatic electromagnetic wave (Volkov wave functions) as the electromagnetic wave is turned on and off. Since the arguments on both sides are delicate, the question for eventual experimental decision is left open. Because the process does have a classical answer, the author would tend to side with those who favor the existence of the IDFS.

VII. REFRACTION OF ELECTRONS BY LASER BEAMS

Up to this point it has been assumed that the laser beam was a linearly polarized monochromatic electromagnetic plane wave. But, in fact, the beam emerging from the end of a laser has a finite beam diameter and, depending on the mode, a complicated transverse structure. An analysis (refs. 46-48) of the equations of motion of an electron in an arbitrary field (not necessarily plane wave) shows that the electron, in addition to its oscillatory motion, responds to gradients of field intensities according to the equation

$$\left\langle \frac{d^2\vec{x}}{dt^2} \right\rangle = -\frac{c^2}{2} \nabla q^2$$

where the brackets average over the purely oscillatory motion and q^2 is the intensity parameter which now varies in different regions of the beam. The equation means that the electron moves, on the average, as if it were in a potential $V = mq^2c^2/2$, all this assuming the averaged motion is non-relativistic. It can be seen that the electron moves <u>away</u> from regions of high intensity.

It can be recognized that the velocity given the electron in the beam direction, as discussed in the previous section, is a special case of this result. For assume that the electron is initially at rest at the point $z = L$ (Figure 4) and that the beam turns on linearly, as it moves, i.e.:

$$q^2 = q_0^2\left(1 - \frac{z}{L}\right) .$$

As the wave overtakes the electron, the electron experiences a constant force in the +z direction of magnitude

$$\nabla V = \frac{1}{2} mc^2 \frac{q_0^2}{L} .$$

The electron, therefore, accelerates at $c^2q_0^2/2L$ cm/sec^2 for a time L/c sec and the velocity reached in this time is $cq_0^2/2$. After time L/c has elapsed, the electron is now in the constant intensity part of the beam, no longer accelerates, and thus drifts constantly at the velocity $cq_0^2/2$. If the beam then passes by and falls to zero intensity in the same way that it built up, the electron will decelerate back to rest. The important point is that any interactions of the electron taking place while it is inside the beam occur in a frame moving with velocity $cq^2/2$ with respect to the lab and this must be taken into account.

A still more interesting effect is that of an electron entering a laser beam from the side. The beam will, for the purposes of this discussion, be modeled as a tube of constant q^2.

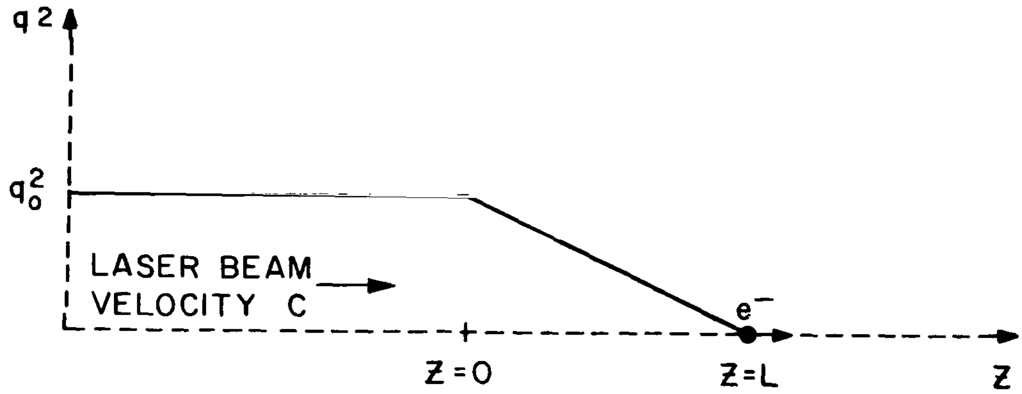

Figure 4

Outside the beam (Figure 5), the electron has kinetic energy

$$T = \frac{1}{2} mv_x^2$$

but inside the kinetic energy is

$$T' = \frac{1}{2} mv_x^2 - \frac{1}{2} mq^2c^2 ,$$

since it has entered the potential (the velocity is lowered by the outward impulsive force the electron receives at the beam surface). It may be noted immediately that, since the kinetic energy can never become negative, if

$$\frac{1}{2} mv_x^2 < \frac{1}{2} mq^2c^2 ,$$

the electron cannot penetrate the beam at all. In other words, electrons slow enough for the condition

$$\frac{v_x^2}{c^2} < q^2$$

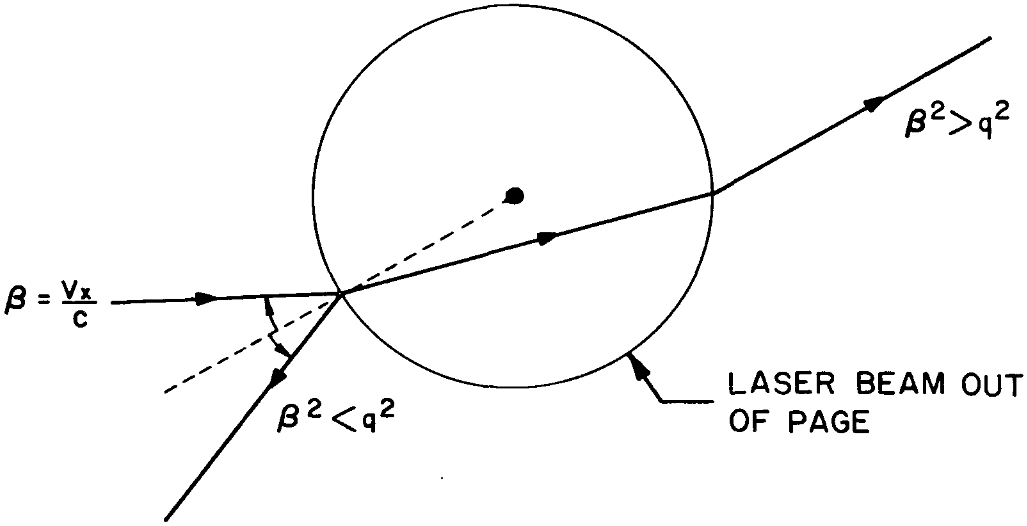

Figure 5

to be true will be totally reflected from the beam. Electrons with

$$\frac{v_x^2}{c^2} > q^2$$

will slow down to velocity

$$\sqrt{v_x^2 - q^2c^2} = v_x \sqrt{\frac{1 - q^2}{\left(\frac{vx}{c}\right)^2}}$$

while inside the beam.

The entire discussion so far is very reminiscent of the inability of light of frequency less than the plasma frequency ω_p to enter an electron gas. The analogy may be drawn tighter and the laser beam treated as a medium of index of refraction

$$\left[1 - \frac{q^2}{\left(\frac{v}{c}\right)^2}\right]^{1/2}$$

with respect to the electron just as an electron gas acts as a medium of index

$$\left[1 - \frac{\omega_p^2}{\omega^2}\right]^{1/2}$$

for a photon. Both ω_p^2 and q^2 depend directly on the density, ω_p^2 on the electron density, and q^2 on the photon density.

All the above considerations can be extended to the region of relativistic electron velocities without any qualitatively new features arising. It may also be noted that the situation directly suggests numerous easy experiments to test the forces on electrons due to gradients in light intensity. Two that immediately come to mind are the scattering of electrons off the gradient produced by focusing a beam, and the reflection of slow electrons from an intense, unfocused beam.

VIII. QUANTUM EFFECTS

Some purely quantum mechanical and quantum field effects which become enhanced or substantially modified in the field of an intense laser beam will be discussed in this section. The considerations below must be considered speculative at best in view of the lack of experimental confirmation of even the classical aspects of high intensity electrodynamics.

Vacuum Polarization

When a bare charge is placed in free space and the electromagnetic interactions turned on, it interacts virtually with all the seething mass of electron-positron pairs, photons, meson pairs, nucleon pairs, and the like, that go into making up what is known as the physical vacuum. The dominant virtual process is the emission and absorption of virtual photons which virtually dissociate into electron-positron pairs (Figure 6).

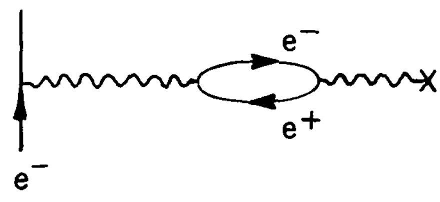

Figure 6

The virtual positrons so produced are attracted to the bare charge and the virtual electrons are repelled; the net effect is the shielding of the bare charge by a small amount. It is this shielded charge that is actually seen as the physical charge of the electron; the bare charge would be seen only if there were no virtual processes going on in the vacuum.

If the electron should now be placed in a very intense laser beam, the rate of virtual photons emitted is stimulated by the laser beam so that more e^+e^- pairs are produced. Thus, the bare electron would be shielded still more and the physical charge of an electron inside a laser beam would be slightly smaller than the physical charge of an electron in free space.

This can be seen quantitatively. In free space, to lowest order, the Fourier transform of the vector potential of the electron is well known (ref. 49):

$$a_{\mu eff}(\nu) = \left(1 - \frac{4\pi}{15} \frac{\alpha\lambda_c^2\nu^2}{c^2}\right) a_\mu(\nu)$$

where $a_\mu(\nu)$ is the Fourier component of the electron's potential at frequency ν in the absence of vacuum polarization, $\alpha = 1/137$ is the fine structure constant, and $\lambda_c = h/mc$ is the Compton wavelength of the electron. If one now simply assumes that the vacuum polarization is stimulated by $n(\nu)\Delta\nu = cI/2h\nu^3$, the total number of laser photons per mode, he finds:

$$a_{\mu\ stim}(\nu) = \left(1 - \frac{4\pi}{15} q^2 \frac{\nu}{\Delta\nu}\right) a_\mu(\nu)$$

where $\Delta\nu$ is the width of the laser spectrum.

Experiments to test this charge screening would be difficult because all the other electron-laser interactions would have to be separated out. Probably the best place to look would be in the changed index of refraction of an electron gas in a laser beam since this index is directly proportional to the electron charge.

Electron-Electron Scattering

Since the coulomb interaction is modified by the presence of an intense field, it has been predicted (refs. 50, 51) that electron-electron (Møller) scattering would be altered by the presence of the laser beam. The interaction under certain conditions could become attractive, and this resonant Møller scattering would become even larger than ordinary Møller scattering.

Pair Production

It has been shown that for $q^2 > 1$ the spectrum of laser light scattered off electrons contains a high energy tail due to higher harmonic production. If the 1,000,000th harmonic is present, real electron-position pairs can be produced in the field of a nucleus. Quantum mechanically, this corresponds to the absorption of 1,000,000 laser photons, each of energy 1 eV, and the subsequent emission of one large 1-MeV photon (Figure 7). This is a species of multiple inverse Compton effect. This effect is particularly interesting because it is a high energy effect, yet it is produced by low energy photons. This is the first case seen so far in which high intensities are converted to high energies and the first indication that high energy physics can be performed by a laser.

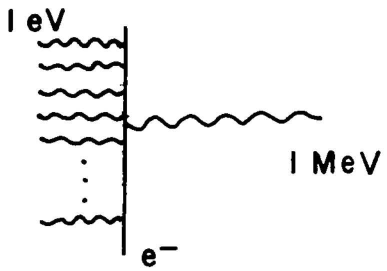

Figure 7

Decay Processes

Because the rate of a decay process is influenced by the final states, and because final states consisting of charged particles are affected by the presence of an intense laser beam, one would expect the decay rates of particles and atoms to be modified. It is predicted that the sensitivity of the rate to the presence of the laser beam increases as the energy available for the decay decreases. It would, therefore, be expected that atomic lifetimes and β-decay lifetimes would be most affected and, in particular, the lifetime of a neutron in an intense laser beam should be lower than that of a neutron in free space.

Photon-Photon Scattering

Photon-photon scattering has long been discussed as one of the cleanest tests of the ideas of quantum electrodynamics because photons mutually interact only through the creation and annihilation of virtual e^+e^- pairs (Figure 8). If the ideas of quantum electrodynamics are wrong, photons would not interact at all. Unfortunately, the cross-section for visible photon-photon scattering is 10^{-65} cm^2, some forty orders of magnitude smaller than that of linear Thomson scattering.

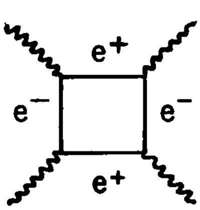

Figure 8

It has been proposed (refs. 52, 53), however, that clashing two intense laser beams in the field of a third intense laser beam would raise the cross-section to where it might become observable. The third beam would stimulate the scattering in its direction so that the probability of scattering becomes proportional to I^3, a very large quantity. The experimental difficulties are clearly formidable.

Bremsstrahlung

The effects of an intense laser beam on bremsstrahlung and inverse bremsstrahlung have been predicted to depend on q^2 through the stimulating effects of the laser beam. Both processes become more probable in an intense beam.

Vacuum Propagation of Light

Finally, worth noting is the possibility that an intense laser beam propagating through the vacuum will diffract less as the intensity is increased. There are two reasons for believing this might be so:

(1) That the beam would increasingly forward-stimulate itself as the intensity is increased, and

(2) That the index of refraction of the vacuum (the vacuum polarization) would increase in the region of highest intensity.

The higher index of refraction would tend to reflect back any of the beam that starts to diffract, the effect being complete when the total internal reflection angle just compensates the diffraction angle. Thus, the same type of circumstance that leads to the well known phenomenon of self-trapping in a liquid could result. At a high enough intensity, the beam would propagate, at a speed slightly less than c, in the vacuum without diffracting.

IX. THE KAPITZA-DIRAC EFFECT

The Kapitza-Dirac effect, or stimulated backward Thomson scattering, occupies a somewhat anomalous position in the field of high intensity electrodynamics and, therefore, its discussion has been left for last. It involves none of the electrodynamical non-linearities previously mentioned and, in fact, occurs for $q^2 \ll 1$. Still, the intensities required are much higher than those available from non-laser sources and therefore the effect might thus be considered an intermediate intensity effect.

If one shoots an electron into the side of a standing laser wave, produced by bouncing a laser off a mirror, the planes of high intensity act as Bragg-diffraction planes and the electron is deflected an angle 2θ (Figure 9), such that

$$\sin\theta = \frac{\lambda_e}{\lambda}$$

where $\lambda_e = h/mv$ is the deBroglie wavelength of the electron and λ is the wavelength of the laser light.

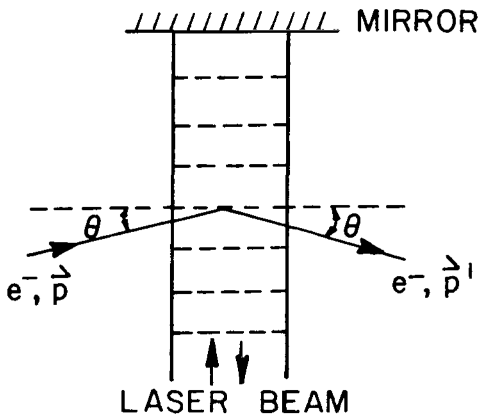

Figure 9

In the photon picture, this deflection comes about because the photons coming from the mirror are Thomson-backscattered off the electron, the probability of such backscattering being increased enormously by the stimulating presence of the photons going toward the mirror. The net change of momentum in each collision is $2h\nu/c$ and the electron must, therefore, recoil by this amount.

Note that the deflection angle is 2θ where

$$\sin\theta = \frac{h}{\lambda p} = \frac{\lambda_e}{\lambda} ,$$

so the scattering looks like first-order Bragg scattering of planes spaced $\lambda/2$ apart (Figure 10).

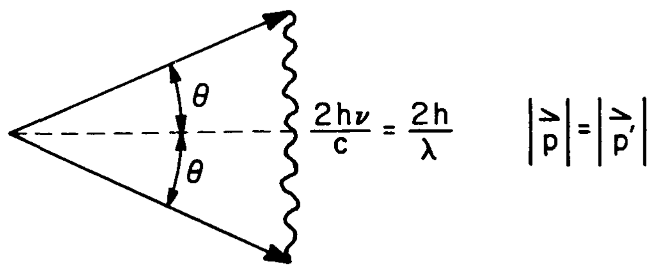

Figure 10

The probability of a single electron being deflected at this angle is given by the probability of backward Thomson scattering times the number of stimulating photons in the backward mode. This latter number is the stimulating factor and is proportional to the intensity. The result for the probability of an electron scattering at the angle θ is (ref. 54):

$$P = \frac{e^4}{2m^2c^2h^2\nu^4} \frac{\ell}{v} \frac{I^2}{\Delta\nu}$$

where

ℓ = the beam diameter

v = the electron velocity

I = the laser beam intensity

$\Delta\nu$ = the laser spectral halfwidth.

P becomes unity for quite moderate beam intensities, namely, $I = 10^8$ W/cm^2 for λ = 1 micron light. The q^2 corresponding to this intensity is only 10^{-10}. For intensities higher than this, P becomes greater than unity and this is to be interpreted as multiple scattering*; the electron will undergo P scatterings as it passes through the beam.

The angle of deflection is only about 0.01 degree for 1 keV electrons and this is what makes the experiment difficult. Thus, despite the moderate power required and the many attempts to perform the experiment, an unambiguous confirmation of the effect has not, to this date, been successful (ref. 56).

More promising approaches to the experiment would be the use of a continuous, high-power CO_2 laser which, when focused, could give the required intensity, or the use of atoms near resonance to increase the cross sections. The smaller angles of deflection for easier-to-handle atoms could be brought to reasonable values by multiple scatterings.

X. CONCLUSION

The aspect of the field of high intensity electrodynamics dealing with the interaction of intense laser beams with free electrons has been reviewed briefly. The interactions of intense beams with bound electrons, i.e., atoms and molecules, which, no doubt, will prove even vaster and practically more important, has not even been considered to this point. For, based on this

*Schappert, to be published.

discussion, it is not difficult to imagine extending these ideas to effects involving level shifts, decay times, stimulated reactions, spin and ionization phenomena, and the like. The effects of intense beams on plasmas, surfaces, semiconductors, and so forth, will certainly be considered after the foundations of the field have set. An understanding of electron and atomic phenomena in intense beams will also be necessary for a microscopic understanding of the burgeoning field of non-linear optics, the theoretical basis of which is now at best phenomenological. Since non-linear optics is already useful in technology and communications, the importance of achieving a microscopic understanding of this field can scarcely be underestimated.

This report will be concluded by again stressing the absence of even a single experimental confirmation of any of the ideas reviewed above. With the development of mode-locked lasers and picosecond pulses, an enormous increase in interest in the field of high intensity electrodynamics can be anticipated, and it is hoped that, within a few years, enough will be understood about the field to make practical applications feasible.

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